

Pressure Stability in Fractional Step Finite Element Methods for Incompressible Flows

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The objective of this paper is to analyze the pressure stability of fractional step finite element methods for incompressible flows that use a pressure Poisson equation. For the classical first-order projection method, it is shown that there is a pressure control which depends on the time step size, and therefore there is a lower bound for this time step for stability reasons. The situation is much worse for a second-order scheme in which part of the pressure gradient is kept in the momentum equation. The pressure stability in this case is extremely weak. To overcome these shortcomings, a stabilized fractional step finite element method is also considered, and its stability is analyzed. Some simple numerical examples are presented to support the theoretical results. © 2001 Academic Press

1. INTRODUCTION

Fractional step methods for the incompressible Navier–Stokes equations have enjoyed widespread popularity since the original works of Chorin [1] and Temam [2]. The reason for this relies on the computational efficiency of these methods (see e.g., [3–6]), basically because of the uncoupling of the pressure from the velocity components. However, several issues related to these methods still deserve further analysis, and perhaps the most salient of these are the behavior of the computed pressure near boundaries and the stability of the pressure itself.

Referring to the pressure boundary conditions, it is well known that numerical boundary layers may appear. The reason for this can be explained by considering the pressure boundary condition associated to the splitting of the continuous problem, as proposed in [1, 2]. Although relevant in some cases, this misbehavior does not affect the global pressure convergence [7], and can be shown to be less dramatic than expected in most situations [8]. In any case, any reference to the correct boundary conditions for the fractional step scheme can be skipped by considering the splitting at the purely algebraic level, once the

space discretization has been performed. This is the approach advocated in [9, 10] and that we will follow here.

The study of the pressure stability for schemes that use a pressure Poisson equation is the main concern of this paper. Surprisingly, this stability is rarely made explicit. It is normally hidden by the convergence analysis, when it is required that the time step size be *small enough*. In general, analyses at the continuous space level are based on the stability of the continuous pressure [11, 12], whereas when the space is discretized it relies on the stability of this discretization, either the finite element method as in [13] or a very simple finite difference setting as in [14]. Other attempts to study the pressure stability are presented in [15], where results much weaker than those presented here are stated (only for the scheme of Section 3.2 and assuming the time step size sufficiently large).

The results to be presented in this paper refer to two types of fractional step schemes, namely, the classical first-order projection method and a second-order algorithm based on the Crank–Nicolson discretization for the viscous and convective terms and a second-order pressure splitting, leaving the pressure gradient at a given time level in the first step and computing its increment in the second one (see [12, 16–19] for different ideas related to second-order schemes). First, the stability of these schemes in the context of a finite element space discretization is analyzed in Section 3 using matrix arguments. It is shown that a certain pressure stability can be expected, regardless of the particular discrete velocity-pressure spaces chosen. In fact, sometimes fractional step schemes that use the pressure Poisson equation are considered as “stabilized” [20]. However, this stability can be useful for the first-order scheme, but it is certainly too weak for the second-order one. In both cases, though, it is given by the time step size, and hence this size is limited from below simply to stabilize the space discretization. It is worth noting that this dependence of the stabilization terms on the time step size also appears in some formulations involving a stabilization of convection, and not the velocity-pressure interpolation (see [21] for a version of the Characteristic-Galerkin method where this fact is clearly demonstrated).

To avoid the bond described, a pressure stabilized scheme is proposed in Section 4. Again, the stability analysis of both the first- and second-order fractional step schemes is undertaken, now using a variational setting rather than the previous matrix language. This stabilization is intended to mimic the stabilizing effect of the first-order projection method. The method was originally presented in [22] for the steady Stokes problem, extended to the nonlinear case in [23], and to the transient problem using a monolithic time discretization in [24]. In spite of the fact that the variational approach used in Section 4 supersedes the matrix analysis of Section 3, the latter is extremely useful to understand the stabilization mechanism introduced by the splitting and, more precisely, by approximation (14) in Section 2.

The final step is to extend the previous stabilization method to the case in which convection needs also to be stabilized. The resulting formulation, based on the ideas of [25], is presented in Section 5. The formulation is able to have control over the pressure gradient and the nonlinear convective term while yielding very accurate numerical results, much less overdiffusive than classical pressure stabilization methods or upwind techniques designed to deal with convection dominated flows.

Some numerical results are presented in Section 6. In spite of their simplicity, they clearly show that the theoretical predictions of the paper are encountered in practice. The summary of the most salient results and the conclusions are finally presented in Section 7.

2. FRACTIONAL STEP METHODS FOR THE NAVIER–STOKES EQUATIONS

2.1. Problem Statement

2.1.1. Continuous problem. In the simplest possible setting, the incompressible Navier–Stokes equations for a fluid moving in a domain Ω of \mathbf{R}^d ($d = 2$ or 3) in a time interval $[0, T]$ are

$$\partial_t \mathbf{u} + \mathbf{u} \cdot \nabla \mathbf{u} - \nu \Delta \mathbf{u} + \nabla p = \mathbf{f} \quad \text{in } \Omega, \quad t \in (0, T), \quad (1)$$

$$\nabla \cdot \mathbf{u} = 0 \quad \text{in } \Omega, \quad t \in (0, T), \quad (2)$$

where \mathbf{u} is the velocity field, p the kinematic pressure, \mathbf{f} the vector field of body forces, and $\nu > 0$ the kinematic viscosity. These equations need to be supplied with an initial condition for the velocity and a boundary condition, which, for simplicity, we will take as the simple homogeneous Dirichlet condition $\mathbf{u} = \mathbf{0}$ on $\partial\Omega$, $t \in [0, T]$.

The finite element space discretization that we will consider is based on the variational formulation of the problem. In order to write it, let us introduce the forms

$$a(\mathbf{u}, \mathbf{v}) := \nu(\nabla \mathbf{u}, \nabla \mathbf{v}), \quad b(q, \mathbf{v}) := (q, \nabla \cdot \mathbf{v}),$$

$$c(\mathbf{u}, \mathbf{v}, \mathbf{w}) := (\mathbf{u} \cdot \nabla \mathbf{v}, \mathbf{w}) + \frac{1}{2}((\nabla \cdot \mathbf{u})\mathbf{v}, \mathbf{w}),$$

where (\cdot, \cdot) denotes the standard L^2 inner product. In these expressions, for a fixed $t \in [0, T]$, $\mathbf{u}, \mathbf{v}, \mathbf{w}$ are assumed to belong to the velocity space $V = [H_0^1(\Omega)]^d$ of vector functions whose components and their derivatives are square-integrable and vanish on $\partial\Omega$, and q belongs to the pressure space $Q = L^2(\Omega)/\mathbf{R}$ of square-integrable functions modulo constants.

Observe that c corresponds to the skew-symmetric form of the convective term. For exactly divergence free \mathbf{u} , the second term of c vanishes. However, it will simplify the analysis of the discrete problem, where we will often make use of the property $c(\mathbf{u}, \mathbf{v}, \mathbf{v}) = 0$ for all $\mathbf{v} \in V$.

Having introduced this notation, the weak form of problem (1)–(2) consists of finding \mathbf{u} and p such that

$$(\partial_t \mathbf{u}, \mathbf{v}) + c(\mathbf{u}, \mathbf{u}, \mathbf{v}) + a(\mathbf{u}, \mathbf{v}) - b(p, \mathbf{v}) = \langle \mathbf{f}, \mathbf{v} \rangle \quad \forall \mathbf{v} \in V,$$

$$b(q, \mathbf{u}) = 0 \quad \forall q \in Q,$$

and \mathbf{u} satisfying the initial condition. The notation $\langle \mathbf{f}, \mathbf{v} \rangle$ stands for the duality pairing of $\mathbf{v} \in V$ and $\mathbf{f} \in [H^{-1}(\Omega)]^d$ (for $t \in [0, T]$).

2.1.2. Monolithic time discretization. We start by considering the simplest time discretization of the problem, namely, the generalized trapezoidal rule, and solving for the velocity and the pressure at the same time. This is what we call *monolithic* scheme. Let $\theta \in [0, 1]$ be a given parameter and consider a partition of $[0, T]$ into N time steps of (for simplicity) equal size δt . Let f be a generic function of time and f^n the value of f at $t^n = n\delta t$ or an approximation to it, and let $f^{n+\theta} := \theta f^{n+1} + (1 - \theta)f^n$, $\delta_t f^n := (f^{n+1} - f^n)/\delta t$. Given \mathbf{u}^n at t^n , the time discrete problem consists of finding \mathbf{u}^{n+1} and p^{n+1} at t^{n+1} as the solution of

$$(\delta_t \mathbf{u}^n, \mathbf{v}) + c(\mathbf{u}^{n+\theta}, \mathbf{u}^{n+\theta}, \mathbf{v}) + a(\mathbf{u}^{n+\theta}, \mathbf{v}) - b(p^{n+1}, \mathbf{v}) = \langle \mathbf{f}^{n+\theta}, \mathbf{v} \rangle \quad \forall \mathbf{v} \in V, \quad (3)$$

$$b(q, \mathbf{u}^{n+1}) = 0 \quad \forall q \in Q. \quad (4)$$

The pressure value computed here has been identified as the pressure evaluated at t^{n+1} , although this is irrelevant for the velocity approximation. The values of interest of θ are $\theta = 1/2$, corresponding to the second-order Crank–Nicolson scheme (see [26] for a thorough analysis of this scheme) and $\theta = 1$, which corresponds to the backward Euler method. If \mathbf{f} is not continuous in time, $\langle \mathbf{f}^{n+\theta}, \mathbf{v} \rangle$ can be taken as the time average of $\langle \mathbf{f}, \mathbf{v} \rangle$ over the time step $[t^n, t^{n+1}]$.

2.1.3. Finite element discretizations. Let \mathbf{V}_h be a finite element space to approximate \mathbf{V} , and Q_h a finite element approximation to Q . Functions in \mathbf{V}_h need to be continuous piecewise polynomials, whereas continuity is in principle not required for Q_h . However, we will consider only continuous pressure interpolations, for reasons explained below.

The finite element discretization of (3) and (4) reads

$$\begin{aligned} (\delta_t \mathbf{u}_h^n, \mathbf{v}_h) + c(\mathbf{u}_h^{n+\theta}, \mathbf{u}_h^{n+\theta}, \mathbf{v}_h) + a(\mathbf{u}_h^{n+\theta}, \mathbf{v}_h) - b(p_h^{n+1}, \mathbf{v}_h) &= \langle \mathbf{f}^{n+1}, \mathbf{v}_h \rangle \quad \forall \mathbf{v}_h \in \mathbf{V}_h, \\ b(q_h, \mathbf{u}_h^{n+1}) &= 0 \quad \forall q_h \in Q_h. \end{aligned}$$

It is well known that for this discrete problem to be stable the velocity and pressure spaces need to satisfy the classical inf-sup condition (see e.g., [27]), which in particular precludes the use of convenient *equal* velocity–pressure interpolations. However, it was early noted that this condition is not required when fractional step methods using a pressure Poisson equation are employed (see e.g., [20, 28–31]). *The analysis and clarification of this situation is precisely the objective of this paper.*

Before presenting the fractional step schemes to be studied, let us introduce the matrix form of the problem. This is given by

$$\mathbf{M} \delta_t \mathbf{U}^n + \mathbf{K}(\mathbf{U}^{n+\theta}) \mathbf{U}^{n+\theta} + \mathbf{G} \mathbf{P}^{n+1} = \mathbf{F}^{n+\theta}, \quad (5)$$

$$\mathbf{D} \mathbf{U}^{n+1} = 0, \quad (6)$$

where \mathbf{U} and \mathbf{P} are the arrays of nodal velocities and pressures, respectively. If we denote the node indexes with superscripts a, b , the space indexes with subscripts i, j , and the standard shape function of node a by N^a , the components of the arrays involved in these equations are

$$\mathbf{M}_{ij}^{ab} = (N^a, N^b) \delta_{ij} \quad (\delta_{ij} \text{ is the Kronecker } \delta),$$

$$\mathbf{K}(\mathbf{U}^{n+\theta})_{ij}^{ab} = (N^a, \mathbf{u}_h^{n+\theta} \cdot \nabla N^b) \delta_{ij} + \frac{1}{2} (N^a, (\nabla \cdot \mathbf{u}_h^{n+\theta}) N^b) \delta_{ij} + \nu (\nabla N^a, \nabla N^b) \delta_{ij},$$

$$\mathbf{G}_i^{ab} = -(\partial_i N^a, N^b),$$

$$\mathbf{F}_i^a = \langle N^a, f_i \rangle,$$

$$\mathbf{D}_j^{ab} = (N^a, \partial_j N^b).$$

It is understood that all the arrays are matrices (except \mathbf{F} , which is a vector) whose components are obtained by grouping together the left indexes in the previous expressions (a and possibly i) and the right indexes (b and possibly j). Note that $\mathbf{G} = -\mathbf{D}^t$.

2.2. Fractional Step Schemes

The fractional step schemes we will consider can be introduced at this point, applied to the fully discrete problem (5)–(6). This is exactly equivalent to

$$M \frac{1}{\delta t} (\hat{U}^{n+1} - U^n) + K(U^{n+\theta})U^{n+\theta} + \gamma GP^n = F^{n+\theta}, \quad (7)$$

$$M \frac{1}{\delta t} (U^{n+1} - \hat{U}^{n+1}) + G(P^{n+1} - \gamma P^n) = 0, \quad (8)$$

$$DU^{n+1} = 0, \quad (9)$$

where \hat{U}^{n+1} is an auxiliary variable and γ is a numerical parameter, whose values of interest are 0 and 1. At this point we can make the essential approximation

$$K(U^{n+\theta})U^{n+\theta} \approx K(\hat{U}^{n+\theta})\hat{U}^{n+\theta}, \quad (10)$$

where $\hat{U}^{n+\theta} := \theta \hat{U}^{n+1} + (1 - \theta)U^n$. Expressing U^{n+1} in terms of \hat{U}^{n+1} using (8) and inserting the result in (9), the set of equations to be solved is

$$M \frac{1}{\delta t} (\hat{U}^{n+1} - U^n) + K(\hat{U}^{n+\theta})\hat{U}^{n+\theta} + \gamma GP^n = F^{n+\theta}, \quad (11)$$

$$\delta t DM^{-1} G(P^{n+1} - \gamma P^n) = D\hat{U}^{n+1}, \quad (12)$$

$$M \frac{1}{\delta t} (U^{n+1} - \hat{U}^{n+1}) + G(P^{n+1} - \gamma P^n) = 0, \quad (13)$$

which have been ordered according to the sequence of solution, for \hat{U}^{n+1} , P^{n+1} , and U^{n+1} . This sequencing of variables has been made possible by (10). The approximation is interpreted in [9, 10] as an *incomplete block LU factorization* of the original problem (7)–(9). The advantage of this discrete approach is that now there is no question about the boundary conditions for the intermediate variable \hat{U}^{n+1} : since boundary conditions are incorporated in the discrete problem (5)–(6), the prescriptions for \hat{U}^{n+1} are exactly the same as for the end-of-step velocity U^{n+1} .

Even though problems (11)–(13) can be implemented as such, it is very convenient to make a further approximation. Observe that $DM^{-1}G$ represents an approximation to the Laplacian operator. In order to avoid dealing with this matrix (which is computationally feasible only if M is approximated by a diagonal matrix), we can approximate

$$DM^{-1}G \approx L, \quad \text{with components } L^{ab} = -(\nabla N^a, \nabla N^b). \quad (14)$$

Matrix L is the standard approximation to the Laplacian operator. The quality of approximation (14) is discussed bellow. Clearly, this approximation is only possible when continuous pressure interpolations are employed. Likewise, it introduces implicitly the same wrong pressure boundary condition as when the splitting is performed at the continuous level (see [3] for a discussion on boundary conditions for the pressure Poisson equation). In [20], the use of approximation (14) is referred to as “approximate projection.”

After using (10) and (14) the problem to be solved is

$$M \frac{1}{\delta t} (\hat{U}^{n+1} - U^n) + K(\hat{U}^{n+\theta})\hat{U}^{n+\theta} + \gamma GP^n = F^{n+\theta}, \quad (15)$$

$$\delta t L(\mathbf{P}^{n+1} - \gamma \mathbf{P}^n) = D \hat{\mathbf{U}}^{n+1}, \quad (16)$$

$$\mathbf{M} \frac{1}{\delta t} (\mathbf{U}^{n+1} - \hat{\mathbf{U}}^{n+1}) + \mathbf{G}(\mathbf{P}^{n+1} - \gamma \mathbf{P}^n) = 0. \quad (17)$$

Once arrived at this point, we may also write the discrete variational equations corresponding to this matrix problem, which are

$$\begin{aligned} & \frac{1}{\delta t} (\hat{\mathbf{u}}_h^{n+1} - \mathbf{u}_h^n, \mathbf{v}_h) + c(\hat{\mathbf{u}}_h^{n+\theta}, \hat{\mathbf{u}}_h^{n+\theta}, \mathbf{v}_h) + a(\hat{\mathbf{u}}_h^{n+\theta}, \mathbf{v}_h) - \gamma b(p_h^n, \mathbf{v}_h) \\ &= (\mathbf{f}^{n+1}, \mathbf{v}_h) \quad \forall \mathbf{v}_h \in \mathbf{V}_h, \\ & -\delta t (\nabla(p_h^{n+1} - \gamma p_h^n), \nabla q_h) = (q_h, \nabla \cdot \hat{\mathbf{u}}_h^{n+1}) \quad \forall q_h \in Q_h, \\ & \frac{1}{\delta t} (\mathbf{u}_h^{n+1} - \hat{\mathbf{u}}_h^{n+1}, \mathbf{v}_h) - b(p_h^{n+1} - \gamma p_h^n, \mathbf{v}_h) = 0 \quad \forall \mathbf{v}_h \in \mathbf{V}_h, \end{aligned}$$

where $\hat{\mathbf{u}}_h$ is obviously the piecewise continuous function obtained by interpolating from the nodal values $\hat{\mathbf{U}}$.

Formally, it is easy to see that the splitting error, introduced by approximation (10), is of order $O(\delta t)$ when $\gamma = 0$ and of order $O(\delta t^2)$ when $\gamma = 1$ (observe from (17) that $O(\|\mathbf{U}^{n+1} - \hat{\mathbf{U}}^{n+1}\|) = \delta t O(\|\mathbf{P}^{n+1} - \gamma \mathbf{P}^n\|)$ in any norm $\|\cdot\|$). In order to have the same error resulting from the splitting and from the original time discretization, we will consider two sets of parameters. The first is $\gamma = 0$, $\theta = 1$, which yields the classical first-order projection scheme discretizing the continuous pressure Poisson equation, and the second $\gamma = 1$, $\theta = 1/2$, which yields a (formally) second-order time accurate method.

3. BASIC STABILITY ESTIMATES

In this section we obtain stability estimates for the two schemes described above. We show that there is a certain pressure stability *regardless of any compatibility requirement between the velocity and pressure approximations*. In order to understand how is this possible, we show first how the fractional step scheme (15)–(17) can be viewed as a *stabilized* monolithic scheme. This equivalence allows us to predict that there will be a pressure stability inherent to the formulations.

3.1. Equivalent Stabilized Monolithic Formulations

In problem (15)–(17) we can eliminate either $\hat{\mathbf{U}}$ or \mathbf{U} , and think of the remaining variable as the approximation to the velocity. These two possibilities lead to two different stabilized formulations, as we show now.

Let us start by writing the problem only in terms of $\hat{\mathbf{U}}$. Since

$$\mathbf{U}^n = \hat{\mathbf{U}}^n - \delta t \mathbf{M}^{-1} \mathbf{G}(\mathbf{P}^n - \gamma \mathbf{P}^{n-1}),$$

Eqs. (15)–(17) can be rewritten as

$$\mathbf{M} \frac{1}{\delta t} (\hat{\mathbf{U}}^{n+1} - \hat{\mathbf{U}}^n) + \mathbf{K}(\hat{\mathbf{U}}^{n+\theta}) \hat{\mathbf{U}}^{n+\theta} + \mathbf{G}[(1 + \gamma) \mathbf{P}^n - \gamma \mathbf{P}^{n-1}] = \mathbf{F}^{n+1}, \quad (18)$$

$$D \hat{\mathbf{U}}^{n+1} - \delta t L(\mathbf{P}^{n+1} - \gamma \mathbf{P}^n) = 0. \quad (19)$$

In the case $\gamma = 0$, this formulation can be viewed as a *stabilized* finite element method, the stabilization effect coming from the pressure Laplacian in the discrete continuity Eq. (19), in a similar way to other popular methods such as Galerkin/Least-Squares [32–34]. Except for the parameter multiplying \mathbf{L} , which now is δt , the formulation is similar to the stabilization method analyzed in [35]. This observation was first pointed out in [36]. Let us remark that the use of $\hat{\mathbf{U}}$ as the velocity variable is also proposed in [37].

An additional comment is that, even though the pressure gradient is treated explicitly in (18), the resulting scheme turns out to be *stable* in time. This results from the forthcoming analysis.

A different approach was advocated in [30]. The idea is to write the problem in terms of \mathbf{U} rather than $\hat{\mathbf{U}}$. Since

$$\hat{\mathbf{U}}^{n+1} = \mathbf{U}^{n+1} + \delta t \mathbf{M}^{-1} \mathbf{G}(\mathbf{P}^{n+1} - \gamma \mathbf{P}^n), \quad (20)$$

Eqs. (15)–(17) can be also rewritten as

$$\begin{aligned} \frac{1}{\delta t}(\mathbf{U}^{n+1} - \mathbf{U}^n) + \mathbf{K}(\mathbf{U}^{n+\theta})\mathbf{U}^{n+\theta} + \mathbf{E}(\mathbf{U}^{n+\theta}) + \mathbf{G}\mathbf{P}^{n+1} &= \mathbf{F}^{n+1}, \\ \mathbf{D}\mathbf{U}^{n+1} + \delta t(\mathbf{D}\mathbf{M}^{-1}\mathbf{G} - \mathbf{L})(\mathbf{P}^{n+1} - \gamma \mathbf{P}^n) &= 0, \end{aligned} \quad (21)$$

where $\mathbf{E}(\mathbf{U}^{n+\theta})$ can be thought of as the splitting error, and is given by

$$\begin{aligned} \mathbf{E}(\mathbf{U}^{n+\theta}) &:= \mathbf{K}(\mathbf{S}^{n+\theta})\mathbf{U}^{n+\theta} + \mathbf{K}(\mathbf{U}^{n+\theta})\mathbf{S}^{n+\theta} + \mathbf{K}(\mathbf{S}^{n+\theta})\mathbf{S}^{n+\theta}, \\ \mathbf{S}^{n+\theta} &:= \theta \delta t \mathbf{M}^{-1} \mathbf{G}(\mathbf{P}^{n+1} - \gamma \mathbf{P}^n). \end{aligned}$$

Clearly, $\mathbf{E}(\mathbf{U}^{n+\theta})$ is formally of order $O(\delta t^{1+\gamma})$.

Consider again the case $\gamma = 0$ (what happens when $\gamma = 1$ is analyzed later). The stabilization effect now comes from the term $\delta t \mathbf{B} \mathbf{P}^{n+1}$, where $\mathbf{B} := \mathbf{D} \mathbf{M}^{-1} \mathbf{G} - \mathbf{L}$. It is shown in [38] that this matrix is *positive semi-definite*. For completeness, let us provide the (simple) proof here. This will allow us to introduce some of the concepts used in Section 4.

Let us consider the vector space $\mathbf{E}_h := \mathbf{V}_h + \nabla \mathbf{Q}_h$, where $\nabla \mathbf{Q}_h$ denotes the space of vector functions which are gradients of functions in \mathbf{Q}_h . If n_1 is the dimension of \mathbf{V}_h and $n_1 + n_{23}$ the dimension of \mathbf{E}_h , it can be split as

$$\mathbf{E}_h = \mathbf{V}_h \oplus \mathbf{V}_h^\perp = \text{Span}\{\mathbf{v}_1, \dots, \mathbf{v}_{n_1}\} \oplus \text{Span}\{\mathbf{v}'_1, \dots, \mathbf{v}'_{n_{23}}\}.$$

Let q_h be an arbitrary element in \mathbf{Q}_h , and \mathbf{Q} the array of nodal values of q_h , and consider the decomposition

$$\nabla q_h = \pi_1 + \pi_{23} = \sum_{k=1}^{n_1} \Pi_{1,k} \mathbf{v}_k + \sum_{k=1}^{n_{23}} \Pi_{23,k} \mathbf{v}'_k, \quad \pi_1 \in \mathbf{V}_h, \quad \pi_{23} \in \mathbf{V}_h^\perp,$$

where Π_1 and Π_{23} are the arrays of nodal values of π_1 and π_{23} , respectively. We have that

$$-\mathbf{Q} \cdot \mathbf{L} \mathbf{Q} = \int_{\Omega} |\nabla q_h|^2 \, d\Omega = \Pi_1 \cdot \mathbf{M} \Pi_1 + \int_{\Omega} \pi_{23} \cdot \pi_{23} \, d\Omega,$$

and, on the other hand, if M_{ij}^{-1} are the components of M^{-1} ,

$$\begin{aligned} Q \cdot DM^{-1}GQ &= - \sum_{i,j=1}^{n_1} \left(\int_{\Omega} \nabla q_h \cdot \mathbf{v}_i \, d\Omega \right) M_{ij}^{-1} \left(\int_{\Omega} \nabla q_h \cdot \mathbf{v}_j \, d\Omega \right) \\ &= - \sum_{i,j=1}^{n_1} \sum_{k,l=1}^{n_1} \Pi_{1,k} \left(\int_{\Omega} \mathbf{v}_k \cdot \mathbf{v}_i \, d\Omega \right) M_{ij}^{-1} \Pi_{1,l} \left(\int_{\Omega} \mathbf{v}_l \cdot \mathbf{v}_j \, d\Omega \right) \\ &= -\Pi_1 \cdot M \Pi_1. \end{aligned}$$

Therefore,

$$Q \cdot BQ = \int_{\Omega} \pi_{23} \cdot \pi_{23} \, d\Omega \geq 0.$$

This proves that B is positive semi-definite and explicitly shows that the components of ∇q_h controlled by this matrix are those orthogonal to the finite element space V_h , that is, π_{23} . This fact is used in [22] to prove that approximation (14) yields an error of the same order as the pressure interpolation error, and thus it does not deteriorate the accuracy of the finite element approximation, although it introduces a spurious pressure behavior near boundaries.

At this point, it is convenient to introduce some additional notation. If X, Y are arrays, $\{X^n\}_{n=0,1,\dots,N}$ is a sequence of arrays of $N+1$ terms and A a symmetric positive semi-definite matrix, we define

$$\|X\|_A := (X \cdot AX)^{1/2},$$

$$\|Y\|_{-A} := \sup_{x \neq 0} \frac{Y \cdot x}{\|x\|_A} \quad (\text{here } A \text{ is assumed to be positive definite}),$$

$$\{X^n\} \in \ell^\infty(A) \Leftrightarrow \|X^n\|_A \leq C < \infty \quad \forall n = 0, 1, \dots, N,$$

$$\{X^n\} \in \ell^p(A) \Leftrightarrow \sum_{n=0}^N \delta t \|X^n\|_A^p \leq C < \infty, \quad 1 \leq p < \infty.$$

Here and in the following, C denotes a positive constant, not necessarily the same at different appearances.

A remark is needed when $A = K$. This matrix is not symmetric, but it has the contribution from the convective term, which is skew-symmetric, and the contribution from the viscous term, K_{visc} , which is symmetric and positive-definite. We will simply write $U \cdot K(U)U = U \cdot K_{\text{visc}}U \equiv \|U\|_K^2$. We will make use also of $L^+ := -L$, which is the positive semi-definite matrix corresponding to the discretization of $-\Delta$.

These definitions will allow us to express our stability results in a compact manner. The basic assumption in all the cases will be that

$$\sum_{n=0}^N \delta t \|F^n\|_{-K}^2 \leq C < \infty, \quad (22)$$

which is the matrix version of the classical condition required for the problem to be well posed.

3.2. Stability of the First-Order Projection Method

We will study now the stability properties of the first-order projection method, which corresponds to scheme (15)–(17) with $\gamma = 0$ and $\delta = 1$:

$$M \frac{1}{\delta t} (\hat{U}^{n+1} - U^n) + K(\hat{U}^{n+1}) \hat{U}^{n+1} = F^{n+1}, \quad (23)$$

$$\delta t L P^{n+1} = D \hat{U}^{n+1}, \quad (24)$$

$$M \frac{1}{\delta t} (U^{n+1} - \hat{U}^{n+1}) + G P^{n+1} = 0. \quad (25)$$

The stability result for this scheme reads:

Stability of the first-order projection scheme:
 $\{U^n\} \in \ell^\infty(M), \quad \{\hat{U}^n\} \in \ell^\infty(M) \cap \ell^2(K), \quad \{\sqrt{\delta t} P^n\} \in \ell^2(L^+)$

The stability estimate for the pressure shows that the pressure gradient *multiplied by δt* is ℓ^2 bounded. When δt is of order $\mathcal{O}(h^2)$ this is optimal [33, 35, 36].

To prove the result, let us multiply (23) by $2\delta t \hat{U}^{n+1}$ and (25) by $2\delta t U^{n+1}$. Using the relation $2a(a-b) = a^2 - b^2 + (a-b)^2$ it is easily obtained that

$$\|\hat{U}^{n+1}\|_M^2 - \|U^n\|_M^2 + \|\hat{U}^{n+1} - U^n\|_M^2 + \delta t \|\hat{U}^{n+1}\|_K^2 \leq \delta t \|F^{n+1}\|_K^2, \quad (26)$$

$$\|U^{n+1}\|_M^2 - \|\hat{U}^{n+1}\|_M^2 + \|U^{n+1} - \hat{U}^{n+1}\|_M^2 + 2\delta t U^{n+1} \cdot G P^{n+1} = 0. \quad (27)$$

On the other hand, from (24) multiplied by P^{n+1} , using the fact that $G = -D^t$, and (20) with $\gamma = 0$, we have

$$\delta t P^{n+1} \cdot L P^{n+1} = -U^{n+1} \cdot G P^{n+1} + \delta t P^{n+1} \cdot D M^{-1} G P^{n+1},$$

$$2\delta t U^{n+1} \cdot G P^{n+1} = 2\delta t^2 P^{n+1} \cdot B P^{n+1} = 2\delta t \|\sqrt{\delta t} P^{n+1}\|_B^2.$$

Using this in (27), adding up (26) and (27) and summing for n (up to any $m \leq N$) it follows that

$$\{U^n\} \in \ell^\infty(M), \quad \{\hat{U}^n\} \in \ell^2(K), \quad \{\sqrt{\delta t} P^n\} \in \ell^2(B).$$

On the other hand, (25) implies

$$\|U^{n+1} - \hat{U}^{n+1}\|_M^2 = -\delta t^2 P^{n+1} \cdot D M^{-1} G P^{n+1},$$

and from the definition of B it is easy to see that

$$\sum_{n=0}^N \delta t \|\sqrt{\delta t} P^{n+1}\|_{L^+}^2 = \sum_{n=0}^N \delta t^2 P^{n+1} \cdot B P^{n+1} + \sum_{n=0}^N \|U^{n+1} - \hat{U}^{n+1}\|_M^2.$$

Both terms on the right-hand side are bounded. The first because we already know that $\{\sqrt{\delta t} P^n\} \in \ell^2(B)$, and the second because it is part of the contribution from (27) when it is added up with (26) and summed for n . Therefore,

$$\{\sqrt{\delta t} P^n\} \in \ell^2(L^+).$$

It remains to show that $\{\hat{U}^n\} \in \ell^\infty(M)$, which follows easily adding up (26) and (27) evaluated at n instead of $n + 1$ and summing for n .

3.3. Stability of the Second-Order Scheme

The second-order scheme corresponds to the choices $\gamma = 1$ and $\theta = 1/2$ in (15)–(17), that is,

$$M \frac{1}{\delta t} (\hat{U}^{n+1} - U^n) + K(\hat{U}^{n+1/2}) \hat{U}^{n+1/2} + GP^n = F^{n+1/2}, \quad (28)$$

$$\delta t L(P^{n+1} - P^n) = D \hat{U}^{n+1}, \quad (29)$$

$$M \frac{1}{\delta t} (U^{n+1} - \hat{U}^{n+1}) + G(P^{n+1} - P^n) = 0. \quad (30)$$

We will take $P^0 = 0$. For this method, the stability result is:

Stability of the second-order scheme:

$$\begin{aligned} \{U^n\} &\in \ell^\infty(M), \quad \{\hat{U}^n\} \in \ell^\infty(M), \quad \{\hat{U}^{n+1/2}\} \in \ell^2(K), \\ \{\delta t P^n\} &\in \ell^\infty(L^+), \quad \{\sqrt{\delta t} \delta P^n\} \in \ell^2(L^+). \end{aligned}$$

Note that the pressure stability is now weaker than in the previous case. Let us prove these estimates. Multiplying (28) by $2\delta t \hat{U}^{n+1/2} = \delta t (\hat{U}^{n+1} + U^n)$ and (30) by $\delta t (U^{n+1} + \hat{U}^{n+1})$ it is found that

$$\begin{aligned} \|\hat{U}^{n+1}\|_M^2 - \|U^n\|_M^2 + \delta t \|\hat{U}^{n+1/2}\|_K^2 + \delta t (\hat{U}^{n+1} + U^n) GP^{n+1} &\leq C \delta t \|F^{n+1}\|_{-K}^2, \\ \|U^{n+1}\|_M - \|U^n\|_M + \delta t (U^{n+1} + \hat{U}^{n+1}) \cdot G(P^{n+1} - P^n) &= 0. \end{aligned} \quad (31)$$

Adding up these two expressions we find that

$$\begin{aligned} \|U^{n+1}\|_M^2 - \|U^n\|_M^2 + \delta t \|\hat{U}^{n+1/2}\|_K^2 + \delta t \hat{U}^{n+1} \cdot GP^{n+1} \\ + \delta t U^n \cdot GP^n + \delta t U^{n+1} \cdot G \delta P^n \leq C \delta t \|F^{n+1}\|_{-K}^2. \end{aligned} \quad (32)$$

From (29) it is easily found that

$$\begin{aligned} \delta t \hat{U}^{n+1} \cdot GP^{n+1} &= -\delta t P^{n+1} \cdot D \hat{U}^{n+1} = -\delta t^2 P^{n+1} \cdot L \delta P^n \\ &= \frac{\delta t^2}{2} (\|P^{n+1}\|_{L^+}^2 - \|P^n\|_{L^+}^2 + \|\delta P^n\|_{L^+}^2). \end{aligned} \quad (33)$$

On the other hand, (30) implies

$$U^{n+1} = \hat{U}^{n+1} - \delta t M^{-1} G \delta P^n,$$

and therefore

$$\begin{aligned} \delta t U^n \cdot GP^n &= -\delta t P^n \cdot D U^n = -\delta t P \cdot D (\hat{U}^n - \delta t M^{-1} G \delta P^{n-1}) \\ &= -\delta t P^n \cdot (\delta t L \delta P^{n-1} - \delta t D M^{-1} G \delta P^{n-1}) \\ &= \delta t^2 P^n \cdot B \delta t P^{n-1} = \frac{\delta t^2}{2} (\|P^n\|_B^2 - \|P^{n-1}\|_B^2 + \|\delta P^{n-1}\|_B^2). \end{aligned} \quad (34)$$

Likewise,

$$\begin{aligned}
 \delta t \mathbf{U}^{n+1} \cdot \mathbf{G} \delta \mathbf{P}^n &= -\delta t \delta \mathbf{P}^n \cdot \mathbf{D} \mathbf{U}^{n+1} = -\delta t \delta \mathbf{P}^n \cdot \mathbf{D} (\hat{\mathbf{U}}^{n+1} - \delta t \mathbf{M}^{-1} \mathbf{G} \delta \mathbf{P}^n) \\
 &= -\delta t \delta \mathbf{P}^n \cdot (\delta t \mathbf{L} \delta \mathbf{P}^n - \delta t \mathbf{D} \mathbf{M}^{-1} \mathbf{G} \delta \mathbf{P}^n) \\
 &= \delta t^2 \delta \mathbf{P}^n \cdot \mathbf{B} \delta \mathbf{P}^n = \delta t \|\sqrt{\delta t} \delta \mathbf{P}^n\|_{\mathbf{B}}^2.
 \end{aligned} \tag{35}$$

Using (33)–(35) in (32), taking $\mathbf{P}^0 = \mathbf{P}^{-1} = 0$ and summing for n we conclude that

$$\{\mathbf{U}^n\} \in \ell^\infty(\mathbf{M}), \quad \{\hat{\mathbf{U}}^{n+1/2}\} \in \ell^2(\mathbf{K}), \quad \{\delta t \mathbf{P}^n\} \in \ell^\infty(\mathbf{L}^+), \quad \{\sqrt{\delta t} \delta \mathbf{P}^n\} \in \ell^2(\mathbf{L}^+).$$

It remains to show that $\{\hat{\mathbf{U}}^n\} \in \ell^\infty(\mathbf{M})$. From (29) we have that

$$\delta t \hat{\mathbf{U}}^{n+1} \cdot \mathbf{G} \delta \mathbf{P}^n = -\delta t \delta \mathbf{P}^n \mathbf{D} \hat{\mathbf{U}}^{n+1} = \delta t^2 \|\delta \mathbf{P}^n\|_{\mathbf{L}^+}^2,$$

which substituted in (31) yields, after making use of (35),

$$\|\hat{\mathbf{U}}^{n+1}\|_{\mathbf{M}}^2 = \|\mathbf{U}^{n+1}\|_{\mathbf{M}}^2 + \delta t^2 \|\delta \mathbf{P}^n\|_{\mathbf{L}^+}^2 + \delta t^2 \|\delta \mathbf{P}^n\|_{\mathbf{B}}^2.$$

The result follows noting that all the right-hand side terms are ℓ^∞ sequences.

4. PRESSURE STABILIZED SCHEMES

4.1. Pressure Stabilization

It has been shown in the previous section that the pressure stability of both the first- and second-order schemes depends on the time step size. It can be anticipated that if δt is very small, stability problems may occur, especially for the second-order scheme. In order to avoid this, one can resort to a *stabilized* formulation, in the same spirit as for monolithic schemes.

Since the stability of the first order scheme is known to be adequate when δt is taken close to the critical time step of the explicit scheme ($\theta = 0$ in (5)), it seems natural to devise a stabilized formulation which inherits the stabilization effect of this first order scheme. This idea was first developed in [22] for the Stokes problem, and is briefly reproduced here for the transient Navier–Stokes equations and used in conjunction with a fractional step method.

Let $\tilde{\mathbf{D}}$, $\tilde{\mathbf{M}}$, and $\tilde{\mathbf{G}}$ be the matrices with the same components as \mathbf{D} , \mathbf{M} , and \mathbf{G} , respectively, but letting the velocity shape functions run over the boundary nodes also. Starting with the monolithic time discretization, the stabilized formulation we consider is

$$\begin{aligned}
 \mathbf{M} \delta_t \mathbf{U}^n + \mathbf{K}(\mathbf{U}^{n+\theta}) \mathbf{U}^{n+\theta} + \mathbf{G} \mathbf{P}^{n+1} &= \mathbf{F}^{n+\theta}, \\
 \mathbf{D} \mathbf{U}^{n+1} + \tau (\tilde{\mathbf{D}} \mathbf{\Pi}^{n+\beta} - \mathbf{L} \mathbf{P}^{n+1}) &= 0, \\
 \tilde{\mathbf{M}} \mathbf{\Pi}^{n+1} &= \tilde{\mathbf{G}} \mathbf{P}^{n+1},
 \end{aligned}$$

where $\mathbf{\Pi}$ is an auxiliary variable, which may be treated either implicitly if $\beta = 1$ or explicitly if $\beta = 0$, and τ is a stabilization parameter which depends on the local element sizes. For the sake of simplicity in the exposition, we will take the finite element meshes quasi-uniform,

and thus τ will depend only on the diameter of the finite element partition h . The stability and convergence analysis dictates that it must behave as

$$\tau \leq C \frac{h^2}{\nu}$$

for viscous dominated cases (an appropriate choice for the constant is $C = 1/4$ when linear elements are used). In general, it can be taken close to the critical time step of the explicit time integration of (5) and (6) (see [22, 24] for further discussion). Nevertheless, the exact value of τ is irrelevant for our discussion.

To see that the modified monolithic problem will inherit the type of stability of the first-order projection method, note that if $\beta = 1$ and Π is eliminated, the modified continuity equation is

$$DU^{n+1} + \tau \tilde{B}P^{n+1} = 0, \quad \tilde{B} := \tilde{D}\tilde{M}^{-1}\tilde{G} - L,$$

which is similar to (21) with $\gamma = 0$.

The reason for having modified matrices D , M and G to \tilde{D} , \tilde{M} and \tilde{G} , respectively, is that in this way the modified monolithic problem does not suffer from spurious pressure boundary layers, as it happens for the first-order projection scheme (see [22] for the analysis of this point).

For the monolithic scheme we may use the same approximations (10) and (14) as for the nonstabilized (Galerkin) formulation to arrive to the split problem:

$$\begin{aligned} M \frac{1}{\delta t} (\hat{U}^{n+1} - U^n) + K(\hat{U}^{n+\theta}) \hat{U}^{n+\theta} + \gamma G P^n &= F^{n+\theta}, \\ \delta t L(P^{n+1} - \gamma P^n) + \tau(LP^{n+1} - \tilde{D}\Pi^{n+\beta}) &= D\hat{U}^{n+1}, \\ M \frac{1}{\delta t} (U^{n+1} - \hat{U}^{n+1}) + G(P^{n+1} - \gamma P^n) &= 0, \\ \tilde{M}\Pi^{n+1} &= \tilde{G}P^{n+1}. \end{aligned}$$

The variational statement corresponding to this matrix problem, which we will use in this section, consists of finding $\hat{u}_h^{n+1} \in V_h$, $p_h^{n+1} \in Q_h$, $u_h^{n+1} \in V_h$, and $\pi_h^{n+1} \in \tilde{V}_h$ such that

$$\begin{aligned} \frac{1}{\delta t} (\hat{u}_h^{n+1} - u_h^n, v_h) + c(\hat{u}_h^{n+\theta}, \hat{u}_h^{n+\theta}, v_h) + a(\hat{u}_h^{n+\theta}, v_h) - \gamma b(p_h^n, v_h) \\ = \langle f^{n+\theta}, v_h \rangle \quad \forall v_h \in V_h, \end{aligned} \quad (36)$$

$$\begin{aligned} -\delta t (\nabla(p_h^{n+1} - \gamma p_h^n), \nabla q_h) - \tau (\nabla p_h^{n+1} - \pi_h^{n+\beta}, \nabla q_h) \\ = (q_h, \nabla \cdot \hat{u}_h^{n+1}) \quad \forall q_h \in Q_h, \end{aligned} \quad (37)$$

$$\frac{1}{\delta t} (u_h^{n+1} - \hat{u}_h^{n+1}, v_h) - b(p_h^{n+1} - \gamma p_h^n, v_h) = 0 \quad \forall v_h \in V_h, \quad (38)$$

$$(\pi_h^{n+1}, \eta_h) = (\nabla p_h^{n+1}, \eta_h) \quad \forall \eta_h \in \tilde{V}_h, \quad (39)$$

where \tilde{V}_h is the space V_h enlarged with the continuous vector functions associated to the boundary nodes. The meaning of the new auxiliary variable π_h is clearly observed from (39): It is the projection of the pressure gradient ∇p_h onto \tilde{V}_h .

In the following analysis, we will need several projections of ∇p_h , all of them with respect to the L^2 -inner product. These projections are

$$\begin{aligned}\pi_1 &: \text{projection onto } V_h, \\ \pi_2 &: \text{projection onto } V_h^\perp \cap \tilde{V}_h, \\ \pi_3 &: \text{projection onto } \tilde{V}_h^\perp.\end{aligned}$$

Likewise, we will denote $\pi_{ij} := \pi_i + \pi_j$. Observe in particular that $\pi_{12} \equiv \pi_h$, and that, using these projections, (38) implies

$$u_h^{n+1} = \hat{u}_h^{n+1} - \delta t(\pi_1^{n+1} - \gamma \pi_1^n). \quad (40)$$

We will see from the following analysis that the terms multiplied by τ provide control on π_3 , that is, the component of the pressure gradient orthogonal to the space of continuous vector fields. Control on π_1 can be directly obtained from the momentum equation, as we will see, whereas control over π_2 follows from the condition

$$\|\nabla p_h\| \leq C(\|\pi_1\| + \|\pi_3\|), \quad (41)$$

which is assumed to hold for any $p_h \in Q_h$. This condition is studied in detail in [22]. In particular, it is shown to hold when equal interpolation is used for the velocity components and the pressure, the situation in which we are interested.

Similarly to Section 3, the stability estimates we will obtain here will be expressed using the following notation. Let $\{f^n\}$ be a sequence of generic functions (scalar or vectorial). We define

$$\begin{aligned}\{f^n\} \in \ell^\infty(L^2) &\Leftrightarrow \|f^n\| \leq C < \infty \quad \forall n = 0, 1, \dots, N, \\ \{f^n\} \in \ell^p(L^2) &\Leftrightarrow \sum_{n=0}^N \delta t \|f^n\|^p \leq C < \infty, \quad 1 \leq p < \infty, \\ \{f^n\} \in \ell^2(H^1) &\Leftrightarrow \sum_{n=0}^N \delta t \|f^n\|_1^2 \leq C < \infty,\end{aligned}$$

where $\|\cdot\|$ is the standard L^2 norm and $\|\cdot\|_1$ is the H^1 norm, that is, the sum of the L^2 norm of a function and the L^2 norm of its gradient. We recall that for functions f vanishing on $\partial\Omega$, $\|f\|_1$ is equivalent to $\|\nabla f\|$. Likewise, we will use the fact that $\ell^2(L^2) \subset \ell^1(L^2)$, valid for finite time intervals of analysis.

Finally, let us state the condition analogous to (22) in the variational setting that we use now, which is

$$\sum_{n=0}^N \delta t \|f^n\|_{-1}^2 \leq C < \infty,$$

where $\|\cdot\|_{-1}$ is the dual norm of $\|\cdot\|_1$.

4.2. Stability of the Stabilized First-Order Scheme

The first-order stabilized scheme is obtained taking $\gamma = 0$ and $\theta = 1$ in (36)–(39), which yields

$$\frac{1}{\delta t}(\hat{\mathbf{u}}_h^{n+1} - \mathbf{u}_h^n, \mathbf{v}_h) + c(\hat{\mathbf{u}}_h^{n+1}, \hat{\mathbf{u}}_h^{n+1}, \mathbf{v}_h) + a(\hat{\mathbf{u}}_h^{n+1}, \mathbf{v}_h) = (\mathbf{f}^{n+1}, \mathbf{v}_h) \quad \forall \mathbf{v}_h \in \mathbf{V}_h, \quad (42)$$

$$-\delta t(\nabla p_h^{n+1}, \nabla q_h) - \tau(\nabla p_h^{n+1} - \boldsymbol{\pi}_h^{n+\beta}, \nabla q_h) = (q_h, \nabla \cdot \hat{\mathbf{u}}_h^{n+1}) \quad \forall q_h \in Q_h, \quad (43)$$

$$\frac{1}{\delta t}(\mathbf{u}_h^{n+1} - \hat{\mathbf{u}}_h^{n+1}, \mathbf{v}_h) - b(p_h^{n+1}, \mathbf{v}_h) = 0 \quad \forall \mathbf{v}_h \in \mathbf{V}_h. \quad (44)$$

The stability result for this formulation is:

Stability of the stabilized first-order scheme:

$$\begin{aligned} \{\mathbf{u}_h^n\} &\in \ell^\infty(L^2), \quad \{\hat{\mathbf{u}}_h^n\} \in \ell^\infty(L^2) \cap \ell^2(H^1), \\ \{\sqrt{\delta t} \nabla p_h^n\} &\in \ell^2(L^2), \quad \{\sqrt{\tau} \nabla p_h^n\} \in \ell^1(L^2) \end{aligned}$$

From this we see that, on top of the stability estimate of the scheme with $\tau = 0$, now we have an additional one which is optimal if $\tau = \mathcal{O}(h^2)$.

Let us proceed to prove these stability results. Taking $\mathbf{v}_h = 2\delta t \hat{\mathbf{u}}_h^{n+1}$ in (42) and $\mathbf{v}_h = 2\delta t \mathbf{u}_h^{n+1}$ in (44) it is found that

$$\|\hat{\mathbf{u}}_h^{n+1}\|^2 - \|\mathbf{u}_h^n\|^2 + \|\hat{\mathbf{u}}_h^{n+1} - \mathbf{u}_h^n\|^2 + \delta t \nu \|\nabla \hat{\mathbf{u}}_h^{n+1}\|^2 \leq C \delta t \|\mathbf{f}^{n+1}\|_{-1}^2, \quad (45)$$

$$\|\mathbf{u}_h^{n+1}\|^2 - \|\hat{\mathbf{u}}_h^{n+1}\|^2 + \|\mathbf{u}_h^{n+1} - \hat{\mathbf{u}}_h^{n+1}\|^2 + \delta t (p_h^{n+1}, \nabla \cdot \mathbf{u}_h^{n+1}) = 0. \quad (46)$$

On the other hand, from (43) and using (40) (with $\gamma = 0$) we have that

$$\begin{aligned} -(p_h^{n+1}, \nabla \cdot \mathbf{u}_h^{n+1}) &= (\nabla p_h^{n+1}, \hat{\mathbf{u}}_h^{n+1}) - \delta t (\nabla p_h^{n+1}, \boldsymbol{\pi}_1^{n+1}) \\ &= \delta t (\nabla p_h^{n+1}, \nabla p_h^{n+1}) + \tau (\nabla p_h^{n+1}, \nabla p_h^{n+1} - \boldsymbol{\pi}_h^{n+\beta}) - \delta t (\boldsymbol{\pi}_1^{n+1}, \boldsymbol{\pi}_1^{n+1}) \\ &= \delta t \|\boldsymbol{\pi}_{23}^{n+1}\|^2 + \tau (\nabla p_h^{n+1}, \nabla p_h^{n+1} - \boldsymbol{\pi}_h^{n+\beta}). \end{aligned} \quad (47)$$

Calling $\bar{\beta} := 1 - \beta$, the second term in the right-hand side can be written as

$$\begin{aligned} (\nabla p_h^{n+1}, \nabla p_h^{n+1} - \boldsymbol{\pi}_h^{n+\beta}) &= (\nabla p_h^{n+1}, \boldsymbol{\pi}_3^{n+1}) + \bar{\beta} (\nabla p_h^{n+1}, \boldsymbol{\pi}_h^{n+1} - \boldsymbol{\pi}_h^n) \\ &= \|\boldsymbol{\pi}_3^{n+1}\|^2 + \frac{\bar{\beta}}{2} (\|\boldsymbol{\pi}_h^{n+1}\|^2 - \|\boldsymbol{\pi}_h^n\|^2 + \|\boldsymbol{\pi}_h^{n+1} - \boldsymbol{\pi}_h^n\|^2). \end{aligned}$$

Using this in (47), the result in (46), adding it with (45), and summing for n it is found that

$$\begin{aligned} \{\mathbf{u}_h^n\} &\in \ell^\infty(L^2), \quad \{\hat{\mathbf{u}}_h^n\} \in \ell^2(H^1), \\ \{\sqrt{\delta t} \boldsymbol{\pi}_{23}^n\} &\in \ell^2(L^2), \quad \{\sqrt{\tau} \boldsymbol{\pi}_3^n\} \in \ell^2(L^2). \end{aligned}$$

On the other hand, from (40) again,

$$\|\mathbf{u}_h^{n+1} - \hat{\mathbf{u}}_h^{n+1}\|^2 = \delta t^2 \|\boldsymbol{\pi}_1^{n+1}\|^2,$$

and since the sum for n of the left-hand side is bounded (it appears when (45) and (46) are added) we have that $\{\sqrt{\delta t}\pi_1^n\} \in \ell^2(L^2)$, which together with the fact that $\{\sqrt{\delta t}\pi_{23}^n\} \in \ell^2(L^2)$, implies that $\{\sqrt{\delta t}\nabla p_h^n\} \in \ell^2(L^2)$.

Stability of $\{\sqrt{\tau}\pi_1^n\}$ is more delicate. It is provided by the momentum equation actually solved by the fractional scheme, which after adding (42) and (44) is found to be

$$\frac{1}{\delta t}(\mathbf{u}_h^{n+1} - \mathbf{u}_h^n, \mathbf{v}_h) + c(\hat{\mathbf{u}}_h^{n+1}, \hat{\mathbf{u}}_h^{n+1}, \mathbf{v}_h) + a(\hat{\mathbf{u}}_h^{n+1}, \mathbf{v}_h) + (\nabla p_h^{n+1}, \mathbf{v}_h) = \langle \mathbf{f}^{n+1}, \mathbf{v}_h \rangle.$$

Taking $\mathbf{v}_h = \pi_1^{n+1}$ and using the standard inverse estimate (see e.g., [39])

$$\|\pi_1^{n+1}\|_1 \leq \frac{C_{\text{inv}}}{h} \|\pi_1^{n+1}\|, \quad (48)$$

valid for quasi-uniform finite element partitions, we have

$$\begin{aligned} \|\pi_1^{n+1}\|^2 &\leq \|\mathbf{f}^{n+1}\|_{-1} \frac{C_{\text{inv}}}{h} \|\pi_1^{n+1}\| + \frac{1}{\delta t} \|\mathbf{u}_h^{n+1} - \mathbf{u}_h^n\| \|\pi_1^{n+1}\| \\ &\quad + (N_a \|\hat{\mathbf{u}}_h^{n+1}\|_1 + N_c \|\hat{\mathbf{u}}_h^{n+1}\|_1^2) \frac{C_{\text{inv}}}{h} \|\pi_1^{n+1}\|, \end{aligned} \quad (49)$$

where N_a and N_c are the norms of the forms a and c , respectively, and we have used the continuity of a and c , that is, $a(\mathbf{u}, \mathbf{v}) \leq N_a \|\mathbf{u}\|_1 \|\mathbf{v}\|_1$, $c(\mathbf{u}, \mathbf{v}, \mathbf{w}) \leq N_c \|\mathbf{u}\|_1 \|\mathbf{v}\|_1 \|\mathbf{w}\|_1$. Dividing by $\|\pi_1^{n+1}\|$, summing for n , noting that $\sqrt{\tau} \leq Ch$ and using the previous stability bounds, it is found that $\{\sqrt{\tau}\pi_1^n\} \in \ell^2(L^2)$. This, together with the fact that $\{\sqrt{\tau}\pi_3^n\} \in \ell^2(L^2)$ and assumption (41), allows us to conclude that $\{\sqrt{\tau}\nabla p_h^n\} \in \ell^2(L^2)$. Observe that the reason why this result can not be improved to $\{\sqrt{\tau}\nabla p_h^n\} \in \ell^2(L^2)$ is the presence of the term $N_c \|\hat{\mathbf{u}}_h^{n+1}\|_1^2$ in (49) (see [24] for a discussion about the possibility to improve this result for the monolithic scheme).

Finally, the fact that $\{\hat{\mathbf{u}}_h^n\} \in \ell^\infty(L^2)$ follows adding (45) and (46) evaluated at n instead of $n+1$ and summing for n .

4.3. Stability of the Stabilized Second-Order Scheme

To conclude this section, let us study the stabilized second-order method, which is obtained selecting $\gamma = 1$ and $\theta = 1/2$ in (36)–(39):

$$\begin{aligned} \frac{1}{\delta t}(\hat{\mathbf{u}}_h^{n+1} - \mathbf{u}_h^n, \mathbf{v}_h) + c(\hat{\mathbf{u}}_h^{n+1/2}, \hat{\mathbf{u}}_h^{n+1/2}, \mathbf{v}_h) + a(\hat{\mathbf{u}}_h^{n+1/2}, \mathbf{v}_h) - b(p_h^n, \mathbf{v}_h) \\ = \langle \mathbf{f}^{n+1/2}, \mathbf{v}_h \rangle \quad \forall \mathbf{v}_h \in \mathbf{V}_h, \end{aligned} \quad (50)$$

$$\begin{aligned} -\delta t(\nabla(p_h^{n+1} - p_h^n), \nabla q_h) - \tau(\nabla p_h^{n+1} - \pi_h^{n+\beta}, \nabla q_h) \\ = (q_h, \nabla \cdot \hat{\mathbf{u}}_h^{n+1}) \quad \forall q_h \in Q_h, \end{aligned} \quad (51)$$

$$\frac{1}{\delta t}(\mathbf{u}_h^{n+1} - \hat{\mathbf{u}}_h^{n+1}, \mathbf{v}_h) - b(p_h^{n+1} - p_h^n, \mathbf{v}_h) = 0 \quad \forall \mathbf{v}_h \in \mathbf{V}_h. \quad (52)$$

Contrary to what happened for the first-order scheme, now *we need to require that* $\tau \leq C\delta t$. The stability we are able to prove in this case is:

Stability of the stabilized second-order scheme:

$$\begin{aligned} \{\mathbf{u}_h^n\} &\in \ell^\infty(L^2), \quad \{\hat{\mathbf{u}}_h^n\} \in \ell^\infty(L^2), \quad \{\hat{\mathbf{u}}_h^{n+1/2}\} \in \ell^2(H^1), \\ \{\delta t \nabla p_h^n\} &\in \ell^\infty(L^2), \quad \{\sqrt{\delta t} \nabla \delta p_h^n\} \in \ell^2(L^2), \quad \{\sqrt{\tau} \nabla p_h^n\} \in \ell^1(L^2) \end{aligned}$$

The useful pressure stability estimate in this case is the last one, which again is optimal when $\tau = \mathcal{O}(h^2)$.

The proof-strategy is similar to the previous cases. Taking $\mathbf{v}_h = 2\delta t \hat{\mathbf{u}}_h^{n+1/2}$ in (50) and $\mathbf{v}_h = \delta t(\mathbf{u}_h^{n+1} + \hat{\mathbf{u}}_h^{n+1})$ in (52) it is found that

$$\begin{aligned} \|\hat{\mathbf{u}}_h^{n+1}\| - \|\mathbf{u}_h^n\| + \delta t \nu \|\nabla \hat{\mathbf{u}}_h^{n+1/2}\|^2 - \delta t(p_h^n, \nabla \cdot \hat{\mathbf{u}}_h^{n+1}) - \delta t(p_h^n, \nabla \cdot \mathbf{u}_h^n) \\ \leq C\delta t \|\mathbf{f}^{n+1/2}\|_{-1}^2, \end{aligned} \quad (53)$$

$$\|\mathbf{u}_h^{n+1}\|^2 - \|\hat{\mathbf{u}}_h^{n+1}\|^2 - \delta t(p_h^{n+1} - p_h^n, \nabla \cdot \mathbf{u}_h^{n+1}) - \delta t(p_h^{n+1} - p_h^n, \nabla \cdot \hat{\mathbf{u}}_h^{n+1}) = 0.$$

Adding these two expressions we get

$$\begin{aligned} \|\mathbf{u}_h^{n+1}\|^2 - \|\mathbf{u}_h^n\|^2 + \delta t \nu \|\nabla \hat{\mathbf{u}}_h^{n+1/2}\|^2 - \delta t(p_h^n, \nabla \cdot \mathbf{u}_h^n) - \delta t(p_h^{n+1} - p_h^n, \nabla \cdot \mathbf{u}_h^{n+1}) \\ - \delta t(p_h^{n+1}, \nabla \cdot \hat{\mathbf{u}}_h^{n+1}) \leq C\delta t \|\mathbf{f}^{n+1/2}\|_{-1}^2. \end{aligned} \quad (54)$$

We need to deal with the different terms involving the pressure in the left-hand side of this inequality. Using (51) and (52), as well as (40), the various terms can be written as

$$\begin{aligned} -(p_h^n, \nabla \cdot \mathbf{u}_h^n) &= -(p_h^n, \nabla \cdot \hat{\mathbf{u}}_h^n) + \delta t(p_h^n, \nabla \cdot (\boldsymbol{\pi}_1^n - \boldsymbol{\pi}_1^{n-1})) \\ &= \delta t(\nabla \delta p_h^{n-1}, \nabla p_h^n) + \tau(\nabla p_h^n - \boldsymbol{\pi}_h^{n+\beta-1}, \nabla p_h^n) \\ &\quad - \delta t(\nabla p_h^n, \boldsymbol{\pi}_1^n - \boldsymbol{\pi}_1^{n-1}) \\ &= \delta t(\boldsymbol{\pi}_{23}^n - \boldsymbol{\pi}_{23}^{n-1}, \nabla p_h^n) + \tau(\nabla p_h^n - \boldsymbol{\pi}_h^{n+\beta-1}, \nabla p_h^n) \\ &= \frac{1}{2}\delta t(\|\boldsymbol{\pi}_{23}^n\|^2 - \|\boldsymbol{\pi}_{23}^{n-1}\|^2 + \|\boldsymbol{\pi}_{23}^n - \boldsymbol{\pi}_{23}^{n-1}\|^2) \\ &\quad + \tau(\nabla p_h^n - \boldsymbol{\pi}_h^{n+\beta-1}, \nabla p_h^n) \end{aligned} \quad (55)$$

$$\begin{aligned} -(\delta p_h^n, \nabla \cdot \mathbf{u}_h^{n+1}) &= -(\delta p_h^n, \nabla \cdot \hat{\mathbf{u}}_h^{n+1}) + \delta t(\delta p_h^n, \nabla \cdot (\boldsymbol{\pi}_1^{n+1} - \boldsymbol{\pi}_1^n)) \\ &= \delta t(\nabla \delta p_h^n, \nabla \delta p_h^n) + \tau(\nabla p_h^{n+1} - \boldsymbol{\pi}_h^{n+\beta}, \nabla \delta p_h^n) \\ &\quad - \delta t(\nabla \delta p_h^n, \boldsymbol{\pi}_1^{n+1} - \boldsymbol{\pi}_1^n) \\ &= \delta t\|\boldsymbol{\pi}_{23}^n - \boldsymbol{\pi}_{23}^{n+1}\|^2 + \tau(\nabla p_h^{n+1} - \boldsymbol{\pi}_h^{n+\beta}, \nabla \delta p_h^n) \end{aligned} \quad (56)$$

$$\begin{aligned} -(p_h^{n+1}, \nabla \cdot \hat{\mathbf{u}}_h^{n+1}) &= \delta t(\nabla \delta p_h^n, \nabla p_h^{n+1}) + \tau(\nabla p_h^{n+1} - \boldsymbol{\pi}_h^{n+\beta}, \nabla p_h^{n+1}) \\ &= \frac{\delta t}{2}(\|\nabla p_h^{n+1}\|^2 - \|\nabla p_h^n\|^2 + \|\nabla p_h^{n+1} - \nabla p_h^n\|^2) \\ &\quad + \tau(\nabla p_h^{n+1} - \boldsymbol{\pi}_h^{n+\beta}, \nabla p_h^{n+1}). \end{aligned} \quad (57)$$

The terms multiplied by τ in these expressions can be written as

$$\begin{aligned}
 (\nabla p_h^{n+1} - \pi_h^{n+\beta}, \nabla p_h^{n+1}) &= (\pi_3^{n+1}, \nabla p_h^{n+1}) + \bar{\beta}(\pi_h^{n+1} - \pi_h^n, \nabla p_h^{n+1}) \\
 &= \|\pi_3^{n+1}\|^2 + \frac{\bar{\beta}}{2}(\|\pi_h^{n+1}\|^2 - \|\pi_h^n\|^2 + \|\pi_h^{n+1} - \pi_h^n\|^2), \\
 (\nabla p_h^{n+1} - \pi_h^{n+\beta}, \nabla \delta p_h^n) &= (\pi_3^{n+1}, \nabla \delta p_h^n) + \bar{\beta}(\pi_h^{n+1} - \pi_h^n, \nabla \delta p_h^n) \\
 &= \frac{1}{2}(\|\pi_3^{n+1}\|^2 - \|\pi_3^n\|^2 + \|\pi_3^{n+1} - \pi_3^n\|^2) + \bar{\beta}\|\pi_h^{n+1} - \pi_h^n\|^2.
 \end{aligned}$$

Using these expressions in (55)–(57), the result in (54), and summing for n it is found that

$$\begin{aligned}
 \{\mathbf{u}_h^n\} &\in \ell^\infty(L^2), \quad \{\hat{\mathbf{u}}_h^{n+1/2}\} \in \ell^2(H^1), \\
 \{\delta t \nabla p_h^n\} &\in \ell^\infty(L^2), \quad \{\sqrt{\delta t} \nabla \delta p_h^n\} \in \ell^2(L^2), \quad \{\sqrt{\tau} \pi_3^n\} \in \ell^2(L^2).
 \end{aligned}$$

On the other hand, from (53), the inequality $ab \leq (a^2 + b^2)/2$ and these stability results it is readily seen that

$$\|\hat{\mathbf{u}}_h^{n+1}\|^2 \leq \frac{1}{2}\|\hat{\mathbf{u}}_h^{n+1}\|^2 + C(\|\mathbf{u}_h^n\|^2 + \delta t^2 \|\nabla p_h^n\|^2 + \delta t \|\mathbf{f}^{n+1/2}\|_{-1}^2),$$

from where $\{\hat{\mathbf{u}}_h^n\} \in \ell^\infty(L^2)$.

It remains to bound π_1^n in $\ell^1(L^2)$. Using the same strategy as for the first-order scheme (see (49)), we now have that

$$\begin{aligned}
 \delta t \sqrt{\tau} \|\pi_1^{n+1}\| &\leq \sqrt{\tau} \|\mathbf{u}_h^{n+1} - \mathbf{u}_h^n\| + \delta t (N_a \|\hat{\mathbf{u}}_h^{n+1/2}\|_1 + N_c \|\hat{\mathbf{u}}_h^{n+1/2}\|_1^2) \\
 &\quad + \delta t \|\mathbf{f}^{n+1/2}\|_{-1}.
 \end{aligned}$$

When summing for n , the only term that still has not been shown to be bounded is the first on the right-hand side. Bounding it will conclude the proof. Since

$$\sum_{n=0}^{N-1} \sqrt{\tau} \|\mathbf{u}_h^{n+1} - \mathbf{u}_h^n\| \leq \sum_{n=0}^{N-1} \sqrt{\tau} \|\mathbf{u}_h^{n+1} - \hat{\mathbf{u}}_h^{n+1}\| + \sum_{n=0}^{N-1} \sqrt{\tau} \|\hat{\mathbf{u}}_h^{n+1} - \mathbf{u}_h^n\| \quad (58)$$

and, from (40) and the assumption $\tau \leq C\delta t$,

$$\begin{aligned}
 \sum_{n=0}^{N-1} \sqrt{\tau} \|\mathbf{u}_h^{n+1} - \hat{\mathbf{u}}_h^{n+1}\| &= \sum_{n=0}^{N-1} \sqrt{\tau} \delta t \|\pi_h^{n+1} - \pi_1^n\| \\
 &\leq C \sum_{n=0}^{N-1} \delta t^2 \|\nabla p_h^{n+1} - \nabla p_h^n\|^2 < \infty,
 \end{aligned}$$

it only remains to bound the last term in (58). Taking $\mathbf{v}_h = \hat{\mathbf{u}}_h^{n+1} - \mathbf{u}_h^n$ in (50) and using an inverse estimate similar to (48), it is found that

$$\begin{aligned}
 \frac{1}{\delta t} \|\hat{\mathbf{u}}_h^{n+1} - \mathbf{u}_h^n\|^2 &\leq \frac{1}{\sqrt{\tau}} A_n \|\hat{\mathbf{u}}_h^{n+1} - \mathbf{u}_h^n\| + B_n, \\
 A_n &:= N_a \|\hat{\mathbf{u}}_h^{n+1/2}\|_1 + N_c \|\hat{\mathbf{u}}_h^{n+1/2}\|_1^2 + \|\mathbf{f}^{n+1/2}\|_{-1}, \\
 B_n &:= (p_h^n, \nabla \cdot \hat{\mathbf{u}}_h^{n+1}) - (p_h^n, \nabla \cdot \mathbf{u}_h^n).
 \end{aligned}$$

Using the fact that if $x^2 \leq bx + c$ then $x \leq 2b + 2\sqrt{c}$, we obtain

$$\sum_{n=0}^{N-1} \sqrt{\tau} \|\hat{\mathbf{u}}_h^{n+1} - \mathbf{u}_h^n\| \leq C \left(\sum_{n=0}^{N-1} \delta t A_n + \sum_{n=0}^{N-1} \sqrt{\delta t} \sqrt{\tau} \sqrt{B_n} \right).$$

The first term on the right-hand side is bounded, and therefore it only remains to obtain a bound for the second. Since $\tau \leq C\delta t$,

$$\sum_{n=0}^{N-1} \sqrt{\delta t} \sqrt{\tau} \sqrt{B_n} \leq C \sum_{n=0}^{N-1} \delta t B_n.$$

From (55) we have

$$\begin{aligned} - \sum_{n=0}^{N-1} \delta t (p_h^n, \nabla \cdot \mathbf{u}_h^n) &\leq \frac{1}{2} \sum_{n=0}^{N-1} \delta t^2 (\|\boldsymbol{\pi}_{23}^n\|^2 - \|\boldsymbol{\pi}_{23}^{n-1}\|^2 + \|\boldsymbol{\pi}_{23}^n - \boldsymbol{\pi}_{23}^{n-1}\|^2) \\ &+ \sum_{n=0}^{N-1} \delta t \tau \left[\|\boldsymbol{\pi}_3^n\|^2 + \frac{\bar{\beta}}{2} (\|\boldsymbol{\pi}_h^n\|^2 - \|\boldsymbol{\pi}_h^{n-1}\|^2 + \|\boldsymbol{\pi}_h^n - \boldsymbol{\pi}_h^{n-1}\|^2) \right] < \infty, \end{aligned}$$

and similarly to (55)–(57) we can obtain

$$\begin{aligned} &\sum_{n=0}^{N-1} \delta t (p_h^n, \nabla \cdot \hat{\mathbf{u}}_h^{n+1}) \\ &= \sum_{n=0}^{N-1} [-\delta t^2 (\nabla \delta p_h^n, \nabla p_h^n) - \delta t \tau (\nabla p_h^{n+1} - \boldsymbol{\pi}_h^{n+\beta}, \nabla p_h^n)] \\ &= \sum_{n=0}^{N-1} \left[\frac{\delta t^2}{2} (\|\nabla p_h^n\|^2 - \|\nabla p_h^{n+1}\|^2 + \|\nabla p_h^n - \nabla p_h^{n+1}\|^2) \right. \\ &\quad \left. - \delta t \tau (\boldsymbol{\pi}_3^{n+1}, \boldsymbol{\pi}_3^n) - \delta t \tau \bar{\beta} (\boldsymbol{\pi}_h^{n+1} - \boldsymbol{\pi}_h^n, \boldsymbol{\pi}_h^n) \right] \\ &\leq C \sum_{n=0}^{N-1} \delta t^2 (\|\nabla p_h^n - \nabla p_h^{n+1}\|^2 + \|\boldsymbol{\pi}_3^{n+1}\|^2 + \|\boldsymbol{\pi}_3^n\|^2 + \bar{\beta} \|\boldsymbol{\pi}_h^{n+1} - \boldsymbol{\pi}_h^n\|^2). \end{aligned}$$

From the results already obtained, the last term on the right-hand side is bounded. This completes the proof of stability.

5. PRESSURE AND CONVECTION STABILIZATION

The pressure stabilization procedure introduced in the previous section consists of adding the term

$$\tau (\nabla p_h^{n+1} - \boldsymbol{\pi}_h^n, \nabla q_h)$$

to the variational equations of the Galerkin method (either to the monolithic or to the fractional step time discretization). We have considered directly the case $\beta = 0$, which is

more appealing from the computational point of view and has the same stability properties as the case $\beta = 1$. This stabilizing term can be thought of as a least-squares form of the component of the pressure gradient orthogonal to the finite element space, $\pi_3^{n+1} = \nabla p_h^{n+1} - \pi_h^{n+1}$. If this allows the stabilization of the pressure, the idea that naturally arises is to use the same strategy to stabilize convection when the convective term dominates the viscous one. This idea is developed in [25]. Here we describe only the final outcome, which consists of adding

$$\tau(\hat{\mathbf{u}}_h^{n+\theta} \cdot \nabla \hat{\mathbf{u}}_h^{n+\theta} + \nabla p_h^{n+1} - \pi_h^n, \hat{\mathbf{u}}_h^{n+\theta} \cdot \nabla \mathbf{v}_h + \nabla q_h)$$

to the variational equations of the Galerkin method. The contribution to the momentum equation is obtained when the pressure test function is $q_h = 0$, whereas the contribution to the continuity equation is found when $\mathbf{v}_h = \mathbf{0}$. The final result is that the equations to be solved when fractional step methods are used are

$$\begin{aligned} & \frac{1}{\delta t}(\hat{\mathbf{u}}_h^{n+1} - \mathbf{u}_h^n, \mathbf{v}_h) + c(\hat{\mathbf{u}}_h^{n+\theta}, \hat{\mathbf{u}}_h^{n+\theta}, \mathbf{v}_h) + a(\hat{\mathbf{u}}_h^{n+\theta}, \mathbf{v}_h) - \gamma b(p_h^n, \mathbf{v}_h) \\ & + \tau(\hat{\mathbf{u}}_h^{n+\theta} \cdot \nabla \hat{\mathbf{u}}_h^{n+\theta} + \nabla p_h^n - \pi_h^n, \hat{\mathbf{u}}_h^{n+\theta} \cdot \nabla \mathbf{v}_h) = \langle \mathbf{f}^{n+\theta}, \mathbf{v}_h \rangle \quad \forall \mathbf{v}_h \in \mathbf{V}_h, \\ & - \delta t(\nabla(p_h^{n+1} - \gamma p_h^n), \nabla q_h) - \tau(\hat{\mathbf{u}}_h^{n+\theta} \cdot \nabla \hat{\mathbf{u}}_h^{n+\theta} + \nabla p_h^{n+1} - \pi_h^n, \nabla q_h) \\ & = (q_h, \nabla \cdot \hat{\mathbf{u}}_h^{n+1}) \quad \forall q_h \in Q_h, \\ & \frac{1}{\delta t}(\mathbf{u}_h^{n+1} - \hat{\mathbf{u}}_h^{n+1}, \mathbf{v}_h) - b(p_h^{n+1} - \gamma p_h^n, \mathbf{v}_h) = 0 \quad \forall \mathbf{v}_h \in \mathbf{V}_h, \\ & (\pi_h^{n+1}, \eta_h) = (\hat{\mathbf{u}}_h^{n+\theta} \cdot \nabla \hat{\mathbf{u}}_h^{n+\theta} + \nabla p_h^{n+1}, \eta_h) \quad \forall \eta_h \in \tilde{\mathbf{V}}_h, \end{aligned}$$

which replace (36)–(39). It has to be observed that the pressure in the first equation is treated explicitly, in order to keep the uncoupling of the velocity and pressure calculations. Observe also that now π_h has the meaning of being the projection of $\hat{\mathbf{u}}_h \cdot \nabla \hat{\mathbf{u}}_h + \nabla p_h$ onto the finite element space.

Let us finally remark that the possibility of treating explicitly the projection π_h greatly simplifies the numerical implementation of this stabilization technique. Apart from the evaluation of the several additional terms in the equations for the velocity and the pressure, only an L^2 projection needs to be performed at the end of each time step. This is computationally very inexpensive, especially if the Gramm matrix $\tilde{\mathbf{M}}$ involved in the process is approximated by the diagonal one obtained from a standard lumping technique.

6. NUMERICAL RESULTS

In this section we present the result of two simple 2-D numerical experiments to illustrate the practical impact of the theoretical findings. The objective of the first is to show the dependence of the pressure stability on the time step, whereas the second is mainly intended to check the accuracy, a point that has been mentioned but not analyzed.

6.1. Cavity Flow Problem

This example is the classical cavity flow problem. The Navier–Stokes equations are solved in the unit square with zero velocities everywhere on the boundary except on the

top, where a tangent unit velocity is prescribed. The viscosity has been taken as $\nu = 0.01$, which yields a Reynolds number of 100. A mesh of 20×20 four-noded bilinear elements has been employed. The flow equations have been advanced in time until the steady-state has been reached. The critical time step for the explicit scheme has been estimated to be

$$\delta t_{\text{crit}} = \left(\frac{4\nu}{h^2} + \frac{2U}{h} \right)^{-1}, \quad (59)$$

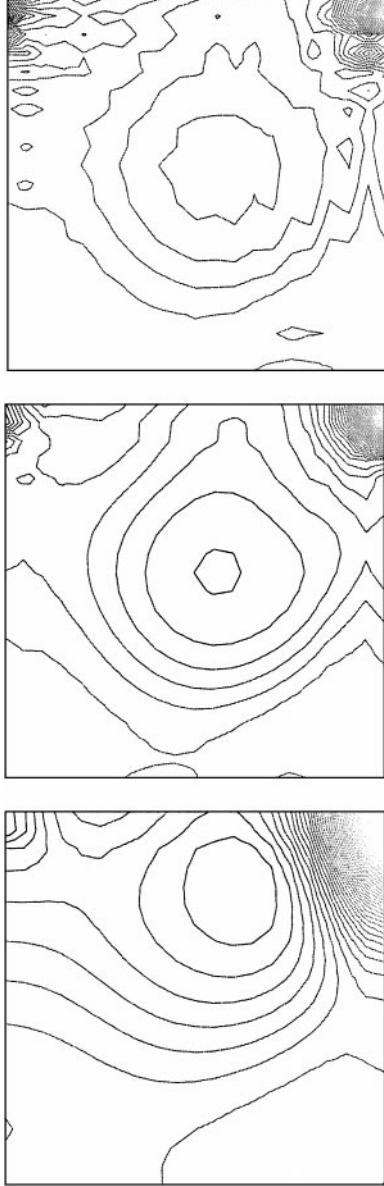


FIG. 1. Pressure contours for the cavity flow-problem using the first-order scheme. From the top to the bottom: $\delta t = 0.1 \delta t_{\text{crit}}$, $\delta t = \delta t_{\text{crit}}$, and $\delta t = 1.0$.

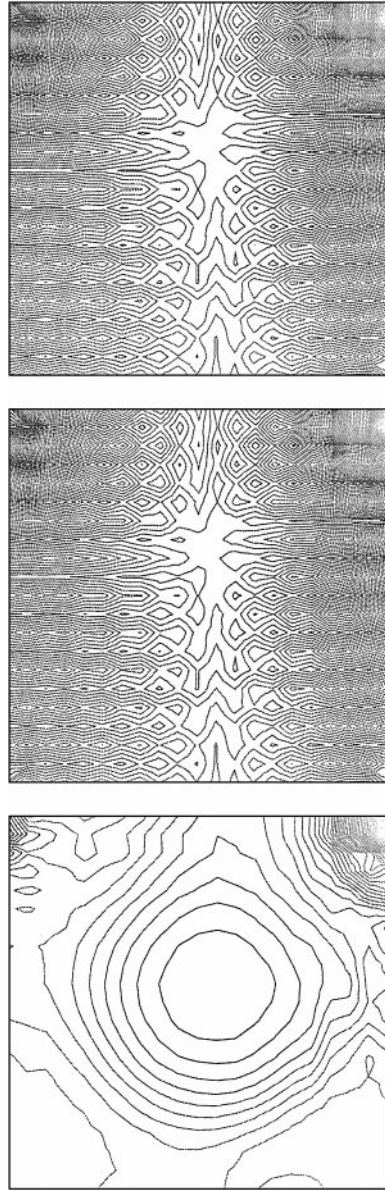


FIG. 2. Pressure contours for the cavity flow problem using the second-order scheme. From the top to the bottom: $\delta t = 0.1 \delta t_{\text{crit}}$, $\delta t = \delta t_{\text{crit}}$, and $\delta t = 1.0$.

where $h = 0.05$ is the element size and $U = 1$ the characteristic velocity. This yields $\delta t_{\text{crit}} = 1/56$.

Pressure contours using the original first-order scheme are shown in Fig. 1. It is clearly observed there that for $\delta t = 0.1 \delta t_{\text{crit}}$ some oscillations appear, whereas for $\delta t = \delta t_{\text{crit}}$ the solution is acceptable. Likewise, when δt is large (1 in this case), the solution is definitely overdiffusive. This shows the dependence of the pressure stability on the time step size.

The same cases have been computed with the second-order scheme. From Fig. 2, it is seen that only for very large values of δt is the weak stability inherent to the scheme activated.

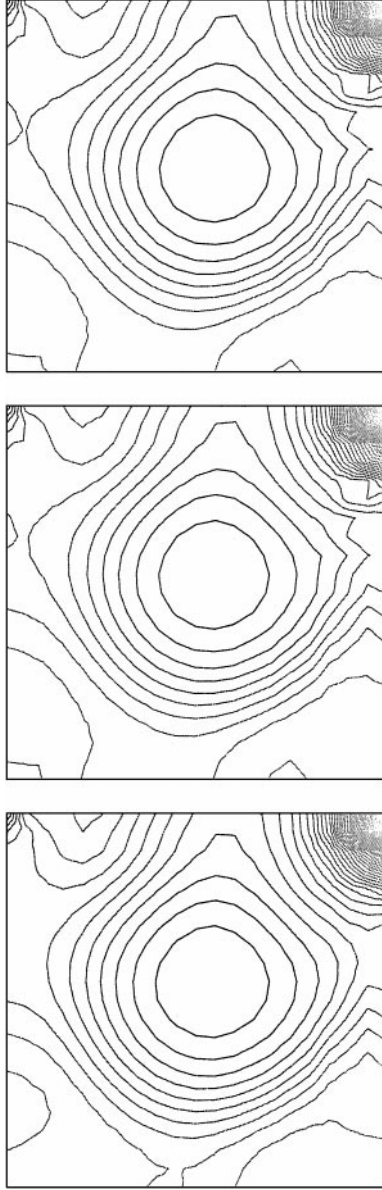


FIG. 3. Pressure contours for the cavity flow problem using the second-order scheme with stabilization. From the top to the bottom: $\delta t = 0.1 \delta t_{\text{crit}}$, $\delta t = \delta t_{\text{crit}}$, and $\delta t = 1.0$.

Contrary to the previous case, for $\delta t = 0.1 \delta t_{\text{crit}}$ and $\delta t = \delta t_{\text{crit}}$ the pressure solution is completely oscillatory.

The final set of results for this problem shown in Fig. 3 corresponds to the stabilized formulation using the second-order scheme. The stability parameter τ has been computed as the critical time step given by (59) but for each element, taking U as the mean element velocity. The pressure solution is correct for all the values of δt .

Even though the implicit schemes we have considered allow us to use time steps as large as desired, it is known that in practice the steady-state is reached faster (both in real time

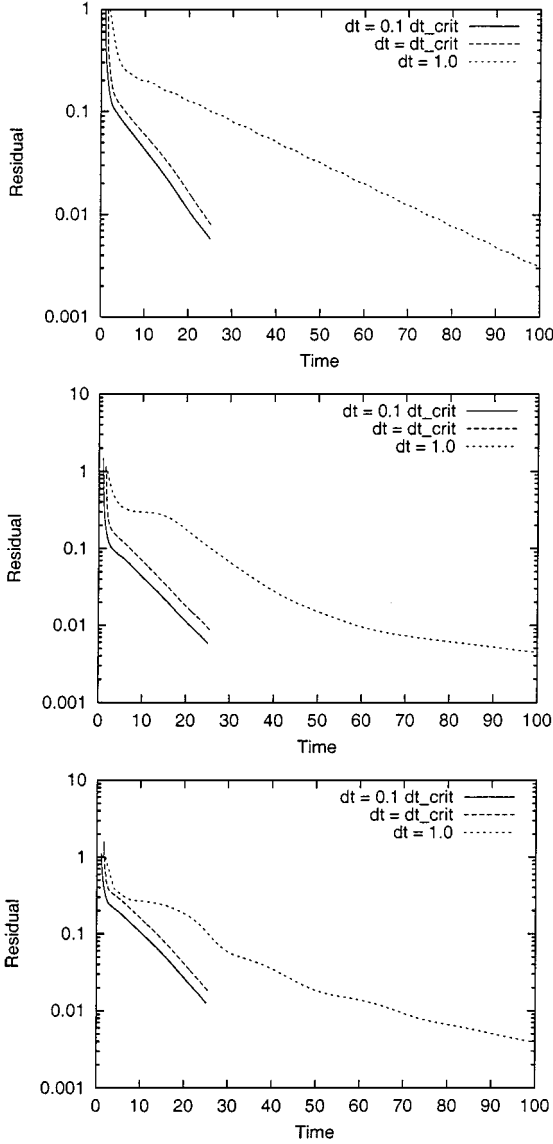


FIG. 4. Velocity convergence toward the steady-state for the cavity flow problem. From the top to the bottom: first-order scheme, second-order scheme, and second-order scheme with stabilization.

and in CPU time) using δt close to δt_{crit} . We have also used this example to demonstrate this. Converge to the steady-state for the different formulations employed is shown in Fig. 4, where the residual is measured as $|\mathbf{U}^n - \mathbf{U}^{n-1}|/|\mathbf{U}^1 - \mathbf{U}^0|$ (here $|\cdot|$ is the standard Euclidean norm of an array). From Fig. 5 it is seen that this dependence on δt is similar for all the schemes.

6.2. A Test with Analytical Solution

We have referred to the two formulations analyzed throughout the paper as the *first-order* projection method and a *second-order* scheme. In this example we test the accuracy of

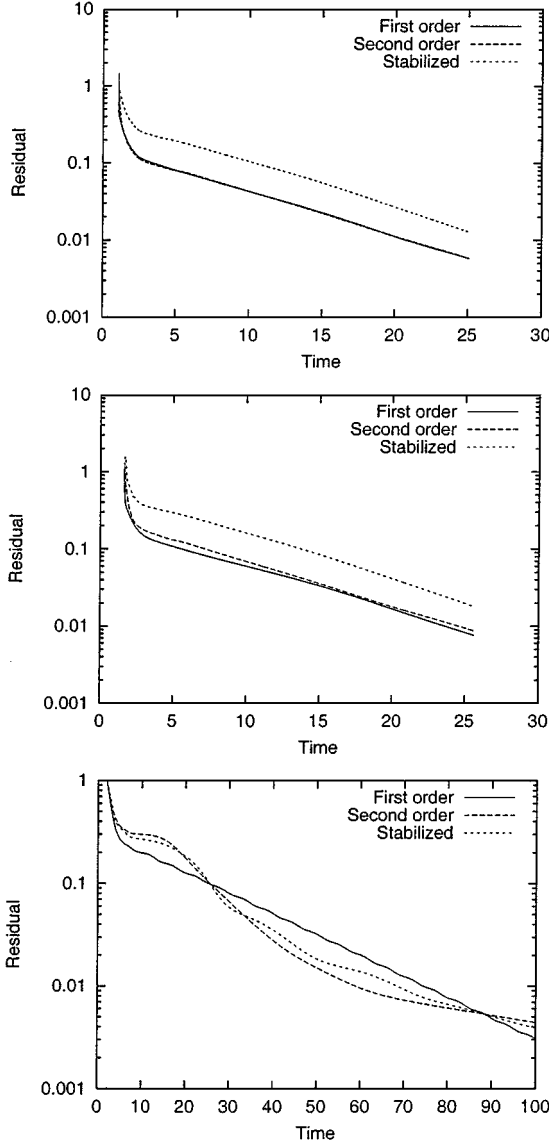


FIG. 5. Velocity convergence toward the steady-state for the cavity flow problem. From the top to the bottom: $\delta t = 0.1 \delta t_{\text{crit}}$, $\delta t = \delta t_{\text{crit}}$, and $\delta t = 1.0$.

these formulations. Since we have seen that the second-order one is unstable, we have combined it with the pressure stabilization technique, with the parameter τ computed as in the previous test.

In this example we solve again the Navier–Stokes equations in the unit square with homogeneous velocity conditions and taking the force vector so as to have as exact solution

$$\begin{aligned}
 u_1(x_1, x_2, t) &= f(x_1)f'(x_2)g(t), & u_2(x_1, x_2, t) &= -f'(x_1)f(x_2)g(t), \\
 f(x) &= 100x^2(1-x)^2, & g(t) &= \cos(\pi t)\exp(-t),
 \end{aligned}
 \tag{60}$$

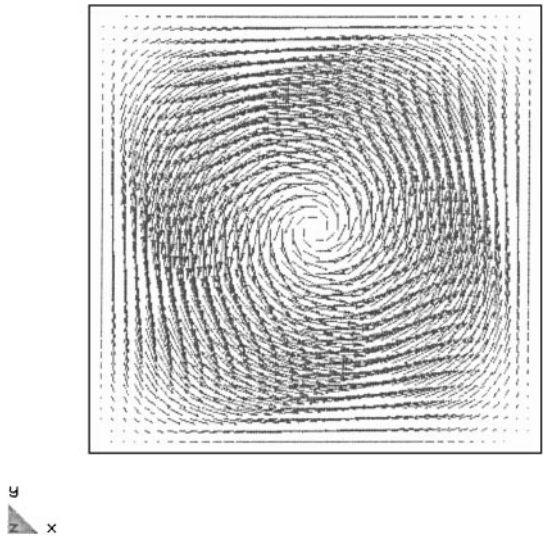


FIG. 6. Velocity vectors for the test with analytical solution.

with two solutions for the pressure, namely, $p = 0$ and $p = 100x^2$. The time interval of analysis is $[0, 1]$ and the viscosity $\nu = 0.1$. A mesh of 40×40 bilinear elements has been employed to discretize the computational domain. The velocity solution at $t = 1$ is shown in Fig. 6.

Convergence of the time approximation for the case $p = 0$ and measured at $t = 1$ is plotted in Fig. 7, where “total” refers to the scheme with $\gamma = 0$ in (15)–(17), and “incremental” refers to the case $\gamma = 1$. The error has been computed as $|\mathbf{U}_{t=1} - \mathbf{U}_{\text{exact}}|/|\mathbf{U}_{\text{exact}}|$, where $\mathbf{U}_{t=1}$ is the numerical solution at $t = 1$, and $\mathbf{U}_{\text{exact}}$ is the array of nodal values of the exact velocity at $t = 1$. Since the exact pressure is $p = 0$, the value of γ does not affect the

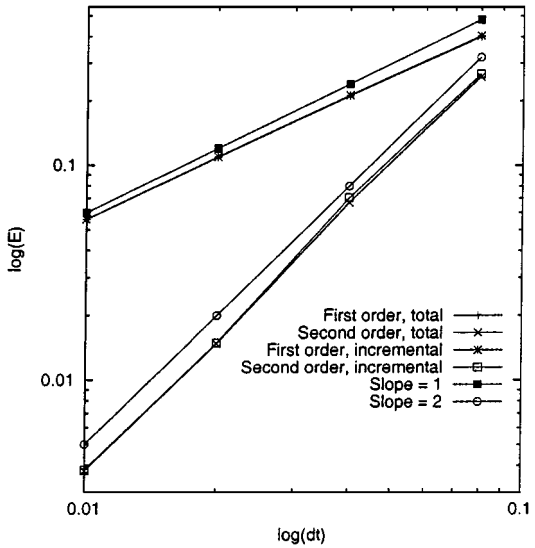


FIG. 7. Convergence for the test with analytical solution. Case $p = 0$.

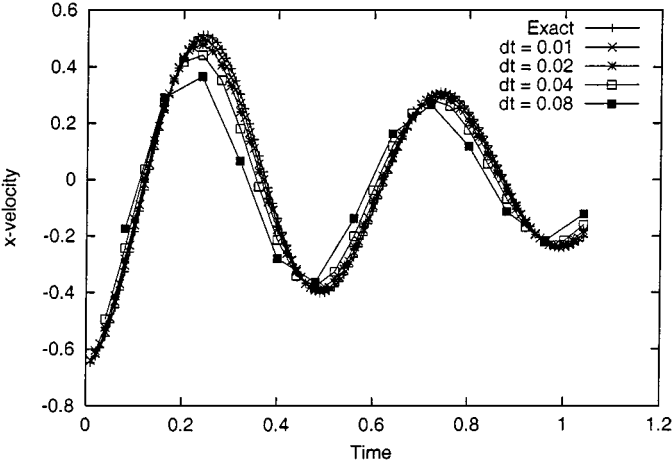


FIG. 8. Evolution of the x -velocity at $(0.75, 0.75)$ using the first-order scheme. Case $p = 0$.

accuracy in this example. From Fig. 7 it is seen that $\theta = 1$ gives a first-order approximation and $\theta = 1/2$ a second-order one. The evolution of the first velocity component u_1 at a point is depicted in Fig. 8 and Fig. 9 for the first- and second-order schemes, respectively.

The evolution to the steady-state when $g(t)$ is replaced by 1 in (60) is shown in Fig. 10. As in the previous example, it is seen that large values of δt may yield slower rates of convergence. Finally, Fig. 11 shows the error at the steady-state as a function of δt when $p = 100x^2$, measuring the well-known property that for the first-order scheme the solution depends largely on δt .

7. SUMMARY OF MAIN RESULTS AND CONCLUSIONS

The main objective of this paper has been to clarify the role of the pressure Poisson equation in the pressure stability of fractional step methods for incompressible flows, in

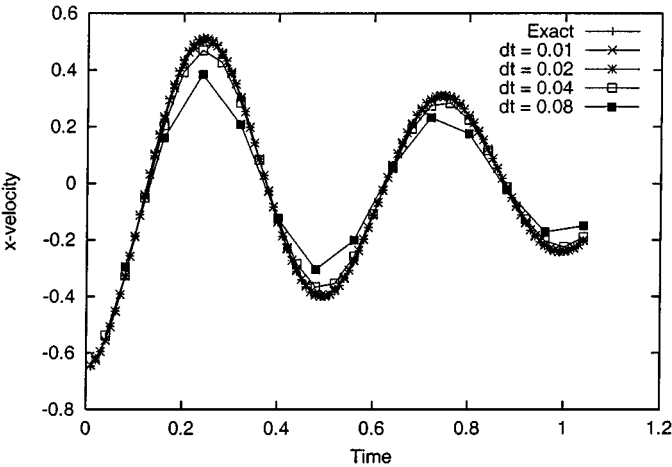


FIG. 9. Evolution of the x -velocity at $(0.75, 0.75)$ using the second-order scheme. Case $p = 0$.

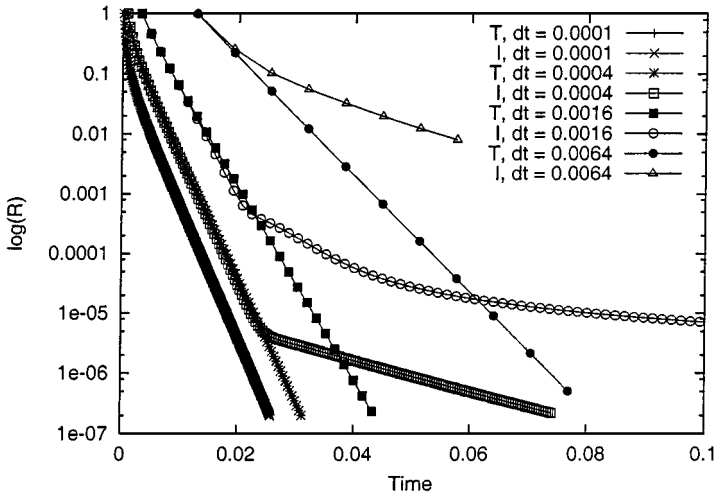


FIG. 10. Convergence to the steady-state for the test with analytical solution (T: total, I: incremental). Case $p = 100x^2$.

our case using the finite element method for the space discretization. The basic stability results for the two schemes analyzed are:

$$\begin{aligned} \text{First order: } & \{\sqrt{\delta t} \nabla p_h^n\} \in \ell^2(L^2), \quad \{\sqrt{\tau} \nabla p_h^n\} \in \ell^1(L^2) \\ \text{Second order: } & \{\delta t \nabla p_h^n\} \in \ell^\infty(L^2), \quad \{\sqrt{\delta t} \nabla \delta p_h^n\} \in \ell^2(L^2), \quad \{\sqrt{\tau} \nabla p_h^n\} \in \ell^1(L^2) \end{aligned}$$

Inspecting these results, the main conclusions that can be drawn include:

- For the original ($\tau = 0$) first-order scheme, pressure is stable, but the parameter that controls the amount of stability is the time step size, and therefore:
- If δt is very small, pressure oscillations may appear.

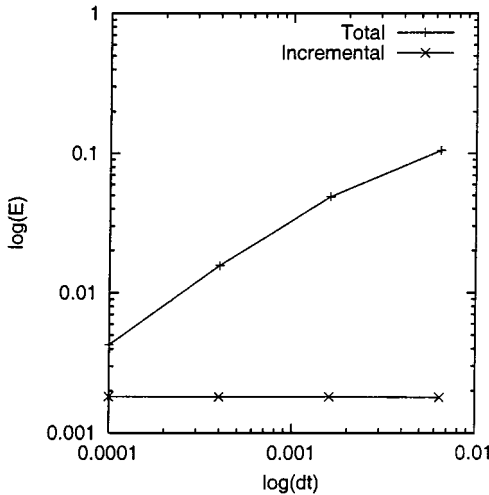


FIG. 11. Error at the steady-state as a function of the time step for the test with analytical solution. Case $p = 100x^2$.

- If δt is large, the method may be overdiffusive. This limits the applicability of implicit schemes, since appropriate values of δt adequate for stability turn out to be close to the critical time step of explicit schemes ($\theta = 0$).

- The original ($\tau = 0$) second-order method has a very poor pressure stability. At the steady-state, only that $\delta t \|\nabla p_h\| \leq C < \infty$ can be ensured, whereas the optimum would be to have control on $\sqrt{\delta t} \|\nabla p_h\|$. However, if δt is very large, the control on $\delta t \|\nabla p_h\|$ can be enough to obtain stable pressures.

- Pressure stability in stabilized schemes depends on an algorithmic parameter τ , which may be chosen independent of δt (except for the condition $\tau \leq C\delta t$ needed for the second-order scheme). Both for the first- and for the second-order methods, this stabilization allows us to free the link stability- δt . For the first-order scheme, δt can be arbitrarily small, and for the second-order one we improve the estimate $\{\delta t \nabla p_h^n\} \in \ell^\infty(L^2)$. In both cases, we are left with $\{\sqrt{\tau} \nabla p_h^n\} \in \ell^1(L^2)$.

A very important fact from the computational point of view is that the pressure gradient projection for stabilized schemes *may be treated explicitly* (which corresponds to taking $\beta = 0$ in (37)). It has been shown that this does not upset stability. Finally, let us mention that convection dominated flows can be stabilized by considering a natural extension of the pressure stabilization technique which has been described in Section 5.

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